

# The Pion Cloud: Insights into Hadron Structure

Anthony W. THOMAS

*Jefferson Lab, 12000 Jefferson Ave., Newport News VA 23606 USA and  
College of William and Mary, Williamsburg VA 23187 USA*

Modern nuclear theory presents a fascinating study in contrasting approaches to the structure of hadrons and nuclei. Nowhere is this more apparent than in the treatment of the pion cloud. As this discussion really begins with Yukawa, it is entirely appropriate that this invited lecture at the Yukawa Institute in Kyoto should deal with the issue.

## §1. Introduction

While the pion has been with us in theory since 1935 and in fact since 1947, chiral symmetry has provided important guidance in dealing with strong interaction phenomena since the 1960s. Well known work by Gell Mann and Lévy, Weinberg, Pagels and many others developed the consequences before the underlying origins in terms of QCD were discovered – for a review see Ref.<sup>1)</sup> That early work was not only mathematically rigorous but was guided by sound physics insight. Since the early 1980s a great deal of work which exploits the chiral symmetry of QCD has been carried out within the framework of chiral perturbation theory. This work is usually described in language which equates it to QCD and talks about a systematic, rigorous description of the hadronic world in terms of a set of low energy constants. In a few outstanding examples this approach has proven very valuable indeed. For some low energy phenomena it has proven possible to make rigorous statements concerning cross sections that had previously seemed incurably model dependent.

Yet, the rise of chiral perturbation theory ( $\chi$ PT) has something of a dark side too. At its best, within a well defined radius of convergence in an expansion in powers of pion mass ( $m_\pi$ ) or momentum transfer ( $Q$ ),  $\chi$ PT provides a mathematically rigorous connection between data sets. Low energy constants derived by fitting one set of data lead to unambiguous predictions elsewhere. In the reductio ad absurdum limit, the low energy world would be reduced to engineering. One would look up a set of low energy constants in the standard reference and turn the handle of a well defined formalism to calculate the result of a new measurement. The low energy constants are by and large without physical content, in many cases scheme or renormalization scale dependent. No physical intuition is either possible or necessary in order to use the formalism. Many physicists find such a world at best dreary or at worst deathly dull. Having as a major aim in life the determination of the as yet unknown (say) 97th low energy constant in a never ending series of physically meaningless parameters does not make getting out of bed in the morning a great joy – for most physicists.

On a more practical level, the systematic order by order expansion involving rigorous application of a strict counting scheme is only meaningful if one lies within

the the corresponding region of convergence. Yet the rigor of going to higher order in a divergent series is rarely discussed. The answer at N<sup>3</sup>LO is not necessarily better than NLO in that case, although it may take several Ph. D. students to get there. The rigor of working consistently within a well defined counting scheme is ephemeral if the counting scheme itself is based on inadequate physics – involving, for example, the wrong degrees of freedom. Classic examples of the latter include studies of spin dependent quantities which define a counting scheme that rigorously excludes the  $\Delta$  resonance.

Of course,  $\chi$ PT itself was in part a natural reaction to the lawlessness of some of the model building that preceded it, where theorists acting more like snake oil salesmen argued the merits of their particular brand in the absence of objective criteria by which to judge. As a particularly extreme reaction to this, some practitioners of  $\chi$ PT go so far as to classify the concept of a pion cloud surrounding hadrons as so ill-defined as to be vigorously eschewed.

In the context of this stimulating meeting at Yukawa's institute it seems worthwhile to re-examine the concept of the pion cloud. In doing so we will find that it has indeed proven its value. Time and again it has predicted new results, later confirmed by experiment or more recently by numerical experiments with QCD itself. It has also provided simple and physically clear and meaningful interpretations of quite diverse phenomena and is still making new predictions which will only be addressed by the next generation of experimental facilities.

## §2. A Framework for Understanding the Role of the Pion Cloud in Hadron Structure

Amongst the many models developed to incorporate the constraints of chiral symmetry into a quantitative picture of hadron structure the cloudy bag model (CBM) was unique.<sup>2)–4)</sup> It did not attempt to describe the short distance structure in terms of pions, since the natural degrees of freedom in QCD are the quarks and gluons. It simply introduced the pion field on top of a successful model of quark confinement (the MIT bag model), in the simplest way consistent with chiral symmetry. It was possible to prove rigorously that, provided the size of the bag (read confinement region) was sufficiently large (say 0.8fm or more), the pion cloud could be treated perturbatively<sup>5)</sup> – the size of the quark core provided a natural mechanism for suppressing the emission of more than one or two pions at a time.

Although the CBM began with the minimal coupling of the pion field to the surface of the bag confining the valence quarks necessary to restore chiral symmetry, it was soon transformed into a form with pseudo-vector coupling throughout the bag volume,<sup>8)</sup> which made it much easier to derive some crucial low energy theorems, such as the Weinberg-Tomozawa formula. (It was also discovered that, at least for massless  $u$  and  $d$  quarks, the results derived using surface coupling were valid in the volume coupled version. For example, the single pion exchange force between two nucleons was identical to that between point-like nucleons for separations greater than twice the bag radius ( $r > 2R$ ), with the form factor arising from finite nucleon size only modifying it when the bags overlapped, in *both* formulations.) The chiral

quark model of Manohar and Georgi,<sup>9)</sup> which was published a few years later, was very similar to that version except that it did not carry out the projection onto colorless, bare hadronic states (the P-space projection).<sup>4)</sup> Only a decade later was it realized<sup>10)</sup> that this projection is essential to ensure the correct infrared limit and hence the correct leading and next-to-leading non-analytic (LNA and NLNA) behaviour of various hadronic properties.<sup>11)</sup> Unfortunately, this lesson is still not widely appreciated and many calculations within so-called chiral quark models are still carried out in a manner which actually violates chiral symmetry!

We now review a large number of important physical results which follow directly from the perturbative nature of the pion cloud found naturally within the cloudy bag model.

### 2.1. *Neutron Charge Distribution*

Within the CBM one had at once a very natural and beautiful way to understand the charge distribution of the neutron.<sup>3)</sup> It had of course been known that the long range negative tail of the neutron was explained in terms of pion emission. But within the Chew-Low model the expansion in terms of the number of pions was divergent and one could say nothing about the short distance structure (inside 1 fm). In the CBM it was suddenly simple; the neutron could be thought of as predominantly a bare neutron bag (with zero charge distribution), but occasionally as a  $\pi^-$  and a bare proton bag whose charge distribution was trivial to calculate. Furthermore, the structure of the pion cloud was such that it peaked in the bag surface, so that the peak of the negative charge distribution of the neutron had a natural interpretation in terms of the bag surface (or the size of the volume within which the valence quarks were confined). On a personal note, I first realized this on New Year's Eve 1980 and had to wait in considerable anticipation for many years before experimental measurements were able to pin down the position of this peak. Modern recoil polarization measurements have established that this peak is indeed around 0.8 fm,<sup>15)</sup> as anticipated in the CBM.

The studies of many different phenomena within the CBM are consistent with a picture in which the total probability of a physical nucleon consisting of a single pion and a bare nucleon is approximately 20% (and a pion and a bare  $\Delta$  around 8-10%).<sup>16)</sup> For the proton (rounding 20% to 21% to simplify the algebra) this means roughly 14%  $\pi^+$ -bare n and 7%  $\pi^0$ -bare p. Similarly for the neutron it's 14%  $\pi^-$ -bare p and 7%  $\pi^0$ -bare n. Taking the charge density at the centre of a bare p to be  $x$  and at the center of a bare n to be 0, this gives for the ratio of the central charge densities of the physical n to p to be approximately  $0.14x/(0.79x + 0.07x) = 0.14/0.86 \sim 1/6$ , which is very close to the experimental ratio extracted from modern studies of the neutron electric form factor at Nikhef, Mainz, MIT-Bates and JLab.<sup>12)-15)</sup> Small corrections may be expected from the  $\pi\Delta$  component of the wave function and from hyperfine effects, which may result in a slightly non-zero charge distribution in the bare n, but the result just derived in this simple and physically transparent manner remains essentially correct.

## 2.2. “Flavor Symmetry” Violations in the Nucleon Sea

Feynman’s much appreciated physical insight led to calculations in the early 1970s of a possible pion contribution to the nucleon sea.<sup>17)</sup> However, this was largely ignored until the discovery of the famous EMC effect generated enormous interest in the possible role of an excess of pions associated with nuclear binding on the structure functions of nuclei.<sup>18),19)</sup> At this time it was realized that the presence of a pion cloud around the nucleon would have profound consequences for the flavor structure of the proton sea. In particular, with the biggest pionic component of the proton wave function being  $\pi^+n$  and the  $\pi^+$  containing only down anti-quarks, it was clear that one expected an excess of anti-down over anti-up quarks in the nucleon sea.<sup>20)</sup> Using the probabilities noted earlier one expected an excess of about 0.14 anti-down quarks.

At the time this novel property of the nucleon sea was predicted it was extremely unusual to think about deep inelastic scattering in terms of hadronic Fock components and certainly not in terms of a pion cloud. It was not until the experimental confirmation of a violation of the Gottfried sum-rule almost 8 years later<sup>22)</sup> that this began to change. Early discussions of the violation of the Gottfried sum-rule often talked about flavor symmetry violation and appeared to confuse that with a violation of isospin, or charge independence. In fact, the simple explanation of how this excess of anti-down quarks was predicted<sup>20),23)</sup> makes it clear that it has *nothing* to do with isospin violation. The 2:1 ratio of the  $\pi^+n$  to  $\pi^0p$  Fock components is a direct consequence of isospin symmetry. In contrast, the term “flavor symmetry” was at best vague and at worst misleading.

In summary, the concept of a pion cloud surrounding the proton led to a simple and natural prediction of an excess of anti-down quarks in the proton sea. The prediction was not only qualitatively but quantitatively in agreement with experiments performed many years later. This discovery has completely changed the standard fits to parton distribution functions.

Perhaps surprisingly, the original paper which predicted the excess of anti-down sea quarks was not primarily concerned with that problem. Its main subject was actually to point out the considerable difference between the strength of the strange and non-strange sea predicted by the meson cloud picture. Indeed, with the hadron size providing a natural high momentum cut-off on the meson-baryon dynamics, it was observed that the strange sea arising from the fluctuation  $p \rightarrow K^+ \Lambda$  would necessarily be much smaller than the non-strange sea associated with the pion cloud.<sup>20)</sup> This was proposed as a natural explanation of the 2:1 ratio that had already been seen in neutrino deep inelastic scattering.<sup>21)</sup> Of course, as well as this sea arising from the meson clouds, which is non-perturbative in the usual sense in which one discusses QCD, there is necessarily a perturbative sea generated through the process  $g \rightarrow q + \bar{q}$  which would be approximately flavor blind (at least for  $u$ ,  $d$  and  $s$ ).

The consequences for the strange sea of the nucleon were explored further by Signal and Thomas just a few years later.<sup>24)</sup> It was noted that while the number of  $s$  and  $\bar{s}$  quarks must be equal in the proton, their momentum distributions will in general not be the same.<sup>25)</sup> This is essentially because of the unequal sharing

of light-cone momentum between the  $K^+$  and  $\Lambda$  in the corresponding piece of the proton wave function. Although we still have no experimental guidance on this issue, it has turned out to be very important in the context of using neutrino deep inelastic scattering to test the Standard Model. It is clearly very important to reliably and accurately determine at least the integral of  $x(s(x) - \bar{s}(x))$  as soon as possible.

### 2.3. Role of the Decuplet in Octet Magnetic Moments

The first systematic inclusion of the decuplet baryons in a calculation of octet magnetic moments was in the study by Th  berge and Thomas.<sup>26)</sup> On physical grounds it was clear that any spin flip quantity would be sensitive to the inclusion of the decuplet. The very convergence of the expansion in pion number for the nucleon was directly related to the role of the  $\Delta$  in the vertex renormalization of the pion-nucleon coupling constant.<sup>4),5)</sup> The results for the octet magnetic moments were in very reasonable agreement with experiment. Within the context of formal  $\chi$ PT it was a decade before decuplet contributions were incorporated in the analysis and even then it took a while to realize that the correct LNA structure required a projection onto bare-baryon pion configurations.

### 2.4. Physically Transparent Estimate of the Strange Quark Content of the Proton

We note first that there is no known example where the current quark masses show up in hadron physics undressed by non-perturbative glue. Thus the cost to make an  $s - \bar{s}$  pair in the proton is of order 1.0 to 1.1 GeV (twice the strange constituent quark mass). On the other hand, creating the  $\bar{s}$  in a kaon and the  $s$  in a  $\Lambda$  costs only 0.65 GeV. (Note that the N to  $K\Sigma$  coupling is considerably smaller than that for N to  $K\Lambda$  and hence, in this simple discussion, we ignore it.) On these grounds alone we expect the virtual transition N to  $K\Lambda$  to dominate the production of strangeness in the proton.

The probability for finding the  $K\Lambda$  configuration is inversely proportional to the excitation energy squared. Naively the transition N to N  $\pi$  costs 140 MeV but with additional kinetic energy this is around 600 MeV in total. Including kinetic energy for the  $K\Lambda$  component as well, it costs roughly twice as much as N  $\pi$ . Thus the  $K\Lambda$  probability is expected to be of order 5%.

We consider first the strangeness radius of the proton, based on this 5%  $K\Lambda$  probability. In the CBM the radius of a  $\Lambda$  bag is about 1 fm, which yields a mean square radius for the strange quark around 0.5 fm<sup>2</sup>. As an estimate of the range of variation possible, we also take the bag radius  $R = 0.8$  fm, with a corresponding mean square radius close to 0.32 fm<sup>2</sup>. In order to estimate the contribution from the kaon cloud, we need to realize that the peak in the Goldstone boson wave function is at the confinement (bag) radius.<sup>6),7)</sup> The meson field then decreases with a range between one over the energy cost of the Fock state and  $1/(m_K + m_\Lambda - m_N)$ . Thus for  $R = 0.8$  fm we get a mean square radius for the  $\bar{s}$  distribution of order 1 fm<sup>2</sup>, while for  $R = 1$  fm we get about 1.4 fm<sup>2</sup>. Weighting the  $s$  by  $-1/3$  and  $\bar{s}$  by  $+1/3$ , we find that the mean square charge radius of strange quarks is between  $(-0.32 + 1.0)/3$  and  $(-0.5 + 1.4)/3$ ; that is, in the range (0.23, 0.30) fm<sup>2</sup>, times the probability for finding the  $K\Lambda$  configuration.

To calculate  $G_E^s$  at  $Q^2 = 0.1 \text{ GeV}^2 = 2.5 \text{ fm}^{-2}$ , we assume that the term  $-Q^2\langle r^2 \rangle/6$  dominates and finally multiply by  $-3$  to agree with the usual convention of removing the strange quark charge. This yields  $G_E^s \in (+0.01, +0.02)$ . It is definitely small and definitely positive for the very clear physical reasons that the  $KA$  probability is small and that the kaon cloud extends outside the  $\Lambda$ . A comparison with the currently preferred fit<sup>27)</sup> to the existing world data<sup>28)–31)</sup> reveals that this estimate has the opposite sign. It is also significantly smaller in magnitude. However, given the current experimental errors, the agreement with data is excellent. (We also note the significant charge symmetry correction in  $^4\text{He}$  reported recently, which would tend to move the central value of the experiment closer to theory.<sup>32)</sup>) Finally, we note that this simple estimate of the mean square strange radius is also in excellent agreement with the recent calculation based on lattice QCD,<sup>33)</sup> namely  $G_E^s(0.1 \text{ GeV}^2) = 0.001 \pm 0.004 \pm 0.004 \text{ fm}^2$ .

Because orbital angular momentum is quantized, the contribution to the magnetic moment from the  $\bar{s}$  in the kaon cloud is much less model dependent. The Clebsch-Gordon coefficients show that in a spin-up proton the probability of a spin down (up)  $\Lambda$ , accompanied by a kaon with orbital angular momentum  $+1$  ( $0$ ), is  $2/3$  ( $1/3$ ). We also know the magnetic moment of the  $\Lambda$  and that it is dominated by the magnetic moment of the  $s$  quark. Hence the total strangeness magnetic moment of the proton is  $-3 \times P_{KA} \times 2/3 \times (+0.6 + 1/3) - 3 \times P_{KA} \times 1/3 \times (-0.6 + 0)$ , where the terms in parentheses are, respectively, the magnetic moment of the spin down (up)  $\Lambda$  and the magnetic moment of the charge  $+1/3$   $\bar{s}$  quark with one unit (or zero units) of orbital angular momentum. The net result, namely  $G_M^s = -0.063 \mu_N$ , is reasonably close to the most recent lattice QCD estimate,<sup>34)</sup> that is,  $G_M^s = -0.046 \pm 0.019 \mu_N$ . From the point of view of this “back of the envelope” estimate, the lattice result clearly has both a natural magnitude and sign.

### §3. Discoveries in Modern Lattice QCD

One of the unexpected but very positive consequences of our *lack* of supercomputing power is the fact that it has not been possible to compute physical hadron properties in lattice QCD. In fact, with computation time scaling like  $m_\pi^{-9}$  (if we include the larger lattice size needed), calculations have covered the pion mass range from  $0.3$  to  $1.0 \text{ GeV}$  (or higher). Far from being a disappointment, this has given us a wealth of unexpected insight into how QCD behaves as the light quark masses are varied.<sup>35)</sup> In terms of the insight this has given us into hadron structure it is both truly invaluable and thus far under-utilized.

The most striking feature of the lattice data is that in the region  $m_\pi > 0.4 \text{ GeV}$ , in fact for almost all of the simulations made so far, all baryon properties show a smooth dependence on quark mass, totally consistent with a constituent quark model. The rapid, non-linear dependence on  $m_\pi$  required by the LNA and NLNA behavior of  $\chi\text{PT}$  are notably absent from the data!

The conventional view of  $\chi\text{PT}$  has no explanation for this simple, universal observation. Worse, in seeking to apply  $\chi\text{PT}$  to extrapolate the data back to the physical pion mass, it has been necessary to rely on ad hoc cancellations between the

high order terms in the usual power series expansion (supplemented by the required non-analytic behavior). In fact, there is strong evidence that such series expansions have been applied well beyond their region of convergence<sup>36)</sup> and that as a result the extrapolations are largely unreliable.

On the other hand, the picture of the pion cloud that we have presented here yields an extremely natural explanation of the universal, constituent quark model behaviour of hadron properties found in the lattice simulations for  $m_\pi > 0.4$  GeV. The natural high momentum cut-off on the momentum of the emitted pion, which is associated with the finite size (typically  $R \sim 1$  fm) of the bare baryon (i.e., the bag in the CBM), strongly suppresses pion loop contributions as  $m_\pi$  increases. The natural mass scale which sets the boundary between rapid chiral variation and constituent quark type behavior is  $1/R \sim 0.2$  to  $0.4$  GeV. Indeed, when in the early investigation of the quark mass dependence of nucleon properties the CBM was compared directly with lattice data, the agreement was remarkably good.<sup>38)</sup> (Similar results have been obtained recently within the chiral quark soliton model.<sup>39)</sup>) We illustrate this in Fig. 1 for the magnetic moments of the proton and neutron. The results were equally as impressive for the N and  $\Delta$  masses and magnetic moments, the proton charge radius and the moments of its parton distribution functions.<sup>40)</sup> The key features necessary to reproduce the behaviour found at large quark mass in lattice QCD *and* to reproduce the experimentally measured data at the physical mass seem to be that:

- The treatment of the pion cloud (chiral) corrections ensures the correct LNA (and NLNA, although in practice this seems less important in many applications) behaviour of QCD
- The pion cloud contribution is suppressed for  $m_\pi$  beyond 0.4 GeV, and
- the underlying quark model exhibits constituent quark like behaviour for the corresponding range of current quark masses.

The CBM satisfies all of these properties.

In practice, of course, in analyzing lattice QCD data one does not want to rely on any particular quark model. However, one does need to suppress the pion cloud

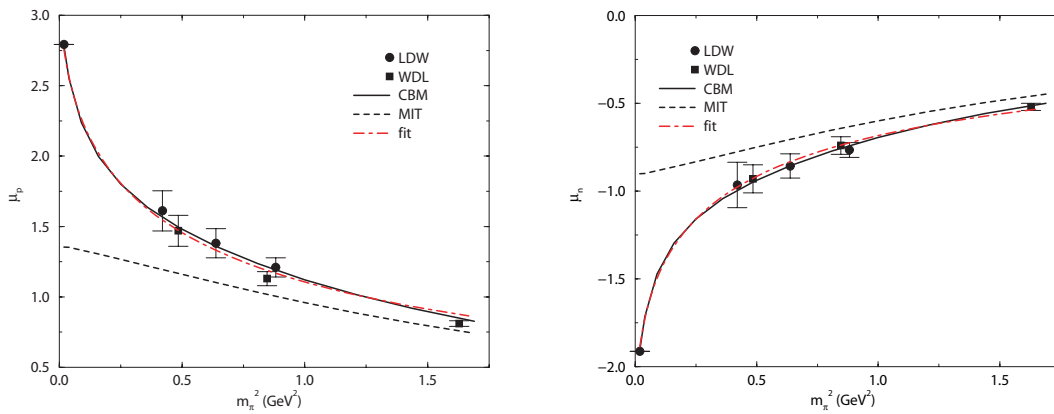


Fig. 1. Comparison between early lattice data for  $\mu_p$  and  $\mu_n$  versus  $m_\pi^2$  and the CBM.<sup>37)</sup> The model is compared with a simple chiral fit and the MIT bag model (without the pion cloud).

as  $m_\pi$  goes up, and the simple use of a finite range regulator (FRR) in the evaluation of the pion loops that yield the LNA and NLNA behaviour ensures this at the cost of one additional parameter, the cut-off mass  $\Lambda$ . If the data are good enough one can use this as a fitting parameter but in general it is sufficient<sup>41)</sup> to choose a value consistent with the physical arguments presented above (e.g.,  $\Lambda$  is  $\sim 0.8$  GeV for a dipole regulator, 0.6 GeV for a monopole and 0.4 GeV using a  $\theta$ -function). The sensitivity of the extrapolation to the choice of the functional form of the FRR is then an additional source of systematic error in the final quoted result. In the case of the nucleon mass the corresponding systematic error was<sup>42)</sup> of the order of a mere 0.1%.

One of the most remarkable results of this physical understanding of the role of the pion cloud and, in particular, its suppression at large pion mass has been the unexpected discovery of a connection between lattice simulations based upon quenched QCD (QQCD) and full QCD.<sup>43)</sup> In a study of the quark mass dependence of the N and  $\Delta$  masses,<sup>44)</sup> it was discovered that if the self-energies appropriate to either QQCD or full QCD were regulated using the same dipole form for the FRR (the dipole being the most natural physical choice given that the axial form factor of the nucleon has a dipole form) with mass parameter  $\Lambda = 0.8$  GeV (the preferred value, as noted above), then the residual expansions for the nucleon mass in QQCD and QCD (and also for the  $\Delta$  in QQCD and QCD) were the same within the errors of the fit! This is a remarkable result which a posteriori gives enormous support to the physical picture of the baryons consisting of confined valence quarks surrounded by a perturbative pion cloud. The baryon core is basically determined by the confinement mechanism and provided the choice of lattice scale reproduces the physically known confining force (either through the string tension or the Sommer parameter,<sup>45)</sup> derived from the heavy quark potential) it makes little difference whether one uses QQCD or full QCD to describe that core. What *does* matter is the change in the chiral coefficients as one goes from QQCD to full QCD.

Perhaps the most significant application of this discovery has been the application to the calculation of the octet magnetic moments and charge radii based on accurate QQCD simulations that extend to rather low quark mass. Using the constraints of charge symmetry this has led to some extremely accurate calculations of the strange quark contributions to the magnetic moment<sup>34)</sup> and charge radius<sup>33)</sup> of the proton. Indeed, those calculations are in excellent agreement with the current world data<sup>27)</sup> but, in a unique example in modern strong interaction physics, they are an order of magnitude more accurate.

#### §4. Concluding Remarks

We have briefly reviewed the power of having a simple and intuitive understanding of hadron structure in which the pion cloud is a crucial element. Time and again the picture has led to new discoveries and predictions. It is extremely unlikely that it will cease to be either useful or inspiring in the near future,<sup>46)</sup> as major new facilities, such as the 12 GeV Upgrade at Jefferson Lab, enable us to probe hadron and nuclear structure in completely new ways.



### Acknowledgements

I would like to thank the many students and colleagues who have collaborated over in the development of the physics described here. I am particularly grateful to Wally Melnitchouk for thoughtful comments on early drafts of this manuscript. This work was supported by the U. S. DOE through contract No. DE-AC-060R23177, under which Jefferson Science Associates operates Jefferson Lab.

### References

- 1) H. Pagels, Phys. Rept. **16**, 219 (1975).
- 2) S. Theberge *et al.*, Phys. Rev. D **22**, 2838 (1980) [Erratum-ibid. D **23**, 2106 (1981)].
- 3) A. W. Thomas, S. Theberge and G. A. Miller, Phys. Rev. D **24** (1981) 216.
- 4) A. W. Thomas, Adv. Nucl. Phys. **13**, 1 (1984).
- 5) L. R. Dodd, A. W. Thomas and R. F. Alvarez-Estrada, Phys. Rev. D **24**, 1961 (1981).
- 6) D. H. Lu, A. W. Thomas and A. G. Williams, Phys. Rev. C **57**, 2628 (1998)
- 7) A. W. Thomas *et al.*, J. Phys. Conf. Ser. **9**, 321 (2005) [arXiv:hep-lat/0502002].
- 8) A. W. Thomas, J. Phys. G **7**, L283 (1981).
- 9) A. Manohar and H. Georgi, Nucl. Phys. B **234**, 189 (1984).
- 10) E. Jenkins and A. V. Manohar, Phys. Lett. B **259**, 353 (1991).
- 11) A. W. Thomas and G. Krein, Phys. Lett. B **456**, 5 (1999) [arXiv:nucl-th/9902013].
- 12) I. Passchier *et al.*, Phys. Rev. Lett. **82**, 4988 (1999) [arXiv:nucl-ex/9907012].
- 13) J. Bermuth *et al.*, Phys. Lett. B **564**, 199 (2003) [arXiv:nucl-ex/0303015].
- 14) H. Zhu *et al.* [E93026 Collaboration], Phys. Rev. Lett. **87**, 081801 (2001)
- 15) J. J. Kelly, AIP Conf. Proc. **698**, 393 (2004).
- 16) J. Speth and A. W. Thomas, Adv. Nucl. Phys. **24**, 83 (1997).
- 17) J. D. Sullivan, Phys. Rev. D **5**, 1732 (1972).
- 18) M. Ericson and A. W. Thomas, Phys. Lett. B **128**, 112 (1983).
- 19) C. H. Llewellyn Smith, Phys. Lett. B **128**, 107 (1983).
- 20) A. W. Thomas, Phys. Lett. B **126**, 97 (1983).
- 21) H. Abramowicz *et al.*, Z. Phys. C **15**, 19 (1982).
- 22) P. Amaudruz *et al.* [New Muon Collaboration], Phys. Rev. Lett. **66**, 2712 (1991).
- 23) A. I. Signal, A. W. Schreiber and A. W. Thomas, Mod. Phys. Lett. A **6**, 271 (1991).
- 24) A. I. Signal and A. W. Thomas, Phys. Lett. B **191**, 205 (1987).
- 25) W. Melnitchouk and M. Malheiro, Phys. Lett. B **451**, 224 (1999) [arXiv:hep-ph/9901321].
- 26) S. Theberge and A. W. Thomas, Nucl. Phys. A **393**, 252 (1983).
- 27) R. D. Young *et al.*, Phys. Rev. Lett. **97**, 102002 (2006) [arXiv:nucl-ex/0604010].
- 28) D. S. Armstrong *et al.* [G0 Collaboration], Phys. Rev. Lett. **95**, 092001 (2005)
- 29) K. A. Aniol *et al.* [HAPPEX Collaboration], Phys. Lett. B **635**, 275 (2006)
- 30) K. A. Aniol *et al.* [HAPPEX Collaboration], Phys. Rev. Lett. **96**, 022003 (2006)
- 31) F. E. Maas *et al.*, Phys. Rev. Lett. **94**, 152001 (2005) [arXiv:nucl-ex/0412030].
- 32) M. Viviani *et al.*, arXiv:nucl-th/0703051.
- 33) D. B. Leinweber *et al.*, Phys. Rev. Lett. **97**, 022001 (2006)
- 34) D. B. Leinweber *et al.*, Phys. Rev. Lett. **94**, 212001 (2005)
- 35) A. W. Thomas, Nucl. Phys. Proc. Suppl. **119**, 50 (2003) [arXiv:hep-lat/0208023].
- 36) D. B. Leinweber, A. W. Thomas and R. D. Young, Nucl. Phys. A **755**, 59 (2005)
- 37) D. B. Leinweber, D. H. Lu and A. W. Thomas, Phys. Rev. D **60**, 034014 (1999)
- 38) D. B. Leinweber *et al.*, Phys. Rev. D **61**, 074502 (2000) [arXiv:hep-lat/9906027].
- 39) K. Goeke *et al.*, Eur. Phys. J. A **27**, 77 (2006); K. Goeke *et al.*, arXiv:hep-ph/0702031.
- 40) W. Detmold, W. Melnitchouk and A. W. Thomas, Phys. Rev. D **66**, 054501 (2002)
- 41) G. P. Lepage, arXiv:nucl-th/9706029.
- 42) D. B. Leinweber, A. W. Thomas and R. D. Young, Phys. Rev. Lett. **92**, 242002 (2004)
- 43) S. Aoki *et al.* [CP-PACS Collaboration], Phys. Rev. Lett. **84**, 238 (2000)
- 44) R. D. Young *et al.*, Phys. Rev. D **66**, 094507 (2002)
- 45) R. Sommer, Nucl. Phys. B **411**, 839 (1994) [arXiv:hep-lat/9310022].
- 46) M. Strikman and C. Weiss, Phys. Rev. D **69**, 054012 (2004) [arXiv:hep-ph/0308191].